

Axion Searches

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The strong CP problem and its resolution through the existence of an axion are briefly reviewed. The constraints on the axion from accelerator searches, from the evolution of red giants and from supernova SN1987a combine to require $m_a < 3 \cdot 10^{-3}$ eV, where m_a is the axion mass. On the other hand, the constraint that axions do not overclose the universe implies $m_a \gtrsim 10^{-6}$ eV. If $m_a \sim 10^{-5}$ eV, axions contribute significantly to the cosmological energy density in the form of cold dark matter. Dark matter axions can be detected by resonant conversion to microwave photons in a cavity permeated by a static magnetic field and tuned to the axion mass. Experiments using this effect are described, as well as several other types of axion searches.

1. Introduction

The axion was postulated nearly two decades ago [1] to explain why the strong interactions conserve P and CP in spite of the fact that the weak interactions violate those symmetries. Consider the Lagrangian of QCD:

$$\mathcal{L}_{QCD} = -\frac{1}{4}G_{\mu\nu}^a G^{a\mu\nu} + \sum_{j=1}^n [\bar{q}_j \gamma^\mu i D_\mu q_j - (m_j q_{Lj}^\dagger q_{Rj} + \text{h.c.})] + \frac{\theta g^2}{32\pi^2} G_{\mu\nu}^a \tilde{G}^{a\mu\nu} \quad . \quad (1)$$

The last term is a 4-divergence and hence does not contribute in perturbation theory. That term does however contribute through non-perturbative effects [2] associated with QCD instantons [3]. Such effects can make the physics of QCD depend upon the value of θ . Using the Adler-Bell-Jackiw anomaly [4], one can show that θ dependence must be present if none of the current quark masses vanishes. Indeed otherwise QCD would have a $U_A(1)$ symmetry and would predict the mass of the η' pseudo-scalar meson to be less than $\sqrt{3}m_\pi \approx 240$ MeV [5], contrary to observation. Using the anomaly, one can further show that QCD depends upon θ only through the combination of parameters:

$$\bar{\theta} = \theta - \arg(m_1, m_2, \dots, m_n) \quad . \quad (2)$$

^{*}This work is supported in part by the U.S. Department of Energy under contract DE-FG05-86ER40272.

If $\bar{\theta} \neq 0$, QCD violates P and CP . The absence of P and CP violations in the strong interactions therefore places an upper limit upon $\bar{\theta}$. The best constraint follows from the experimental bound [6] on the neutron electric dipole moment which yields: $\bar{\theta} < 10^{-9}$.

The question then is: why is $\bar{\theta}$ so small? In the Standard Model of particle interactions, the quark masses originate in the electroweak sector which violates P and CP . There is no reason why the overall phase of the quark mass matrix should exactly match the value of θ from the QCD sector to yield $\bar{\theta} < 10^{-9}$. In particular, if CP violation is introduced in the manner of Kobayashi and Maskawa [7], the Yukawa couplings that give masses to the quarks are arbitrary complex numbers and hence $\arg \det m_q$ and $\bar{\theta}$ are expected to be of order one.

The problem why $\bar{\theta} < 10^{-9}$ is usually referred to as the “strong CP problem”. The existence of an axion solves this problem in a simple manner which is rich in implications for experiment, for astrophysics and for cosmology. There are two alternative solutions however. The first alternative is to set $m_u = 0$. This removes the θ -dependence of QCD and thus solves the strong CP problem. The well-known calculation of the pseudo-scalar meson masses in lowest order of chiral perturbation theory yields $m_u \simeq 4$ MeV, which is incompatible with $m_u = 0$. This calculation also predicts the successful Gell-Mann - Okubo relation

among the pseudo-scalar masses squared. It is possible to have $m_u = 0$ by invoking second order effects [8]. This is a reasonable proposition because m_s happens to be of order the QCD scale. However, when second order effects are included [9], the Gell-Mann - Okubo relation is in general violated. Thus the price for having $m_u = 0$ through higher order effects is that the Gell-Mann - Okubo relation becomes an accident. The second alternative solution to the strong CP problem is to assume that CP and/or P is spontaneously broken but is otherwise a good symmetry. In this case, $\bar{\theta}$ is calculable and may be arranged to be small [10].

In addition, it is worth emphasizing that the strong CP problem need not be solved in the low energy theory. Indeed, as Ellis and Gaillard [11] pointed out, if in the standard model $\bar{\theta} = 0$ near the Planck scale then $\bar{\theta} \ll 10^{-9}$ at the QCD scale.

Peccei and Quinn [12] proposed to solve the strong CP problem by postulating the existence of a global $U_{PQ}(1)$ quasi-symmetry. $U_{PQ}(1)$ must be a symmetry of the theory at the classical (i.e., at the Lagrangian) level, it must be broken explicitly by those non-perturbative effects that make the physics of QCD depend upon θ , and finally it must be spontaneously broken. The axion [13] is the quasi-Nambu-Goldstone boson associated with the spontaneous breakdown of $U_{PQ}(1)$. One can show that, if a $U_{PQ}(1)$ quasi-symmetry is present, then

$$\bar{\theta} = \theta - \arg(m_1 \dots m_n) - \frac{a(x)}{f_a}, \quad (3)$$

where $a(x)$ is the axion field and $f_a = v/N$ is called the axion decay constant. v is the vacuum expectation value which spontaneously breaks $U_{PQ}(1)$ and N is an integer which expresses the color anomaly of $U_{PQ}(1)$. Axion models have N degenerate vacua [14]. The non-perturbative effects that make QCD depend upon $\bar{\theta}$ produce an effective potential $V(\bar{\theta})$ whose minimum is at $\bar{\theta} = 0$. [15] Thus, by postulating an axion, $\bar{\theta}$ is allowed to relax to zero dynamically and the strong CP problem is solved.

The properties of the axion can be derived using the methods of current algebra [16]. The ax-

ion mass is given in terms of f_a by

$$m_a \simeq 0.6 \text{ eV} \frac{10^7 \text{ GeV}}{f_a}. \quad (4)$$

All the axion couplings are inversely proportional to f_a . Of particular interest here is the axion coupling to two photons:

$$\mathcal{L}_{a\gamma\gamma} = -g_\gamma \frac{\alpha a(x)}{\pi f_a} \vec{E} \cdot \vec{B} \quad (5)$$

where \vec{E} and \vec{B} are the electric and magnetic fields, α is the fine structure constant, and g_γ is a model-dependent coefficient of order one. $g_\gamma = 0.36$ in the DFSZ model [17] whereas $g_\gamma = -0.97$ in the KSVZ model [18]. The coupling of the axion to a spin 1/2 fermion f has the form:

$$\mathcal{L}_{a\bar{f}f} = ig_f \frac{m_f}{v} a \bar{f} \gamma_5 f \quad (6)$$

where g_f is a model-dependent coefficient of order one. In the KSVZ model the coupling to electrons is zero at tree level. Models with this property are called 'hadronic'.

2. Constraints from laboratory searches and astrophysics

The searches for the axion in high energy and nuclear physics experiments are discussed in the reviews by Kim [1] and by Peccei [1]. A complete list of experiments can be found in the Review of Particle Physics [19]. If the axion is heavier than 1 MeV and decays quickly into e^+e^- (lifetime of order 10^{-11} sec or less), then it is ruled out by negative searches for rare particle decays such as $\pi^+ \rightarrow a(e^+e^-)e^+\nu_e$ [20]. The rate of this reaction follows simply from the mixing of a and π^0 and the known decay $\pi^+ \rightarrow \pi^0 e^+\nu_e$, and hence it is almost completely model-independent [21]. Alternatively, the axion is long lived (10^{-11} sec or more). In this case it is severely constrained by negative searches in beam dumps. In a beam dump, axions are produced by many different processes involving independent couplings, such as $a - \pi^0$, $a - \eta$ and $a - \eta'$ mixing, the axion couplings to two photons and to two gluons, and the couplings to quarks and gluons. It is difficult to calculate these processes with precision at all energies. On the other hand, the many processes

add up incoherently and they can not all vanish. Thus one can give a reliable estimate for the total production ($p+N \rightarrow a+X$, or $e+N \rightarrow a+X$) and interaction ($a+N \rightarrow X$) cross-sections. Many such searches [19] have been carried out. They rule out axions with mass larger than about 50 keV.

The astrophysical constraints on the axion are described in detail in the reviews by Turner [1] and Raffelt [1]. Axions are emitted by stars in a variety of processes such as Compton-like scattering ($\gamma + e \rightarrow a + e$), axion bremsstrahlung ($e + N \rightarrow N + e + a$) and the Primakoff process ($\gamma + N \rightarrow N + a$). The rate at which a star burns its nuclear fuel is limited by the rate at which it can lose the energy produced. Emission of light weakly coupled bosons, such as the axion, allows a star to radiate energy efficiently because such particles can escape the star at once (without rescattering) whereas photons are emitted only from the stellar surface [22]. Thus the existence of an axion accelerates stellar evolution, which may be inconsistent with observation. The longevity of red giants rules out the mass range $200 \text{ keV} \gtrsim m_a \gtrsim 0.5 \text{ eV}$ [23,24] for hadronic axions. Above 200 keV the axion is too heavy to be copiously emitted in the thermal processes taking in red giants, whereas below 0.5 eV it is too weakly coupled. For axions with a large coupling to electrons [$g_e = 0(1)$ in Eq. 6] the ruled out range can be extended to $200 \text{ keV} \gtrsim m_a \gtrsim 10^{-2} \text{ eV}$ because axion emission through the Compton-like process $\gamma + e \rightarrow a + e$ cools the helium core to such an extent as to prevent the onset of helium burning [25].

Finally the range $2 \text{ eV} \gtrsim m_a \gtrsim 3 \cdot 10^{-3} \text{ eV}$ is ruled out by Supernova 1987a [26]. The constraint follows from the fact that the duration of the associated neutrino events in the large underground proton decay detectors [27] is consistent with theoretical expectations based on the premise that the collapsed supernova core cools by emission of neutrinos. If the axion mass is in the above-mentioned range, the core cools instead by axion emission and the neutrino burst is excessively shortened. The supernova constraint is quite axion model-independent because the axions are emitted by axion bremsstrahlung in

nucleon-nucleon scattering ($N+N \rightarrow N+N+a$) and the relevant couplings follow simply from the mixing of the axion with the π^0 which is a general feature of axion models.

When the limits from laboratory searches are combined with the astrophysical constraints, all of the axion mass range down to approximately $3 \cdot 10^{-3} \text{ eV}$ is ruled out.

3. Axion cosmology

The implications of the existence of an axion for the history of the early universe may be briefly described as follows. At a temperature of order v , a phase transition occurs in which the $U_{PQ}(1)$ symmetry becomes spontaneously broken. This is called the PQ phase transition. At these temperatures, the non-perturbative QCD effects which produce the effective potential $V(\theta)$ are suppressed [28], the axion is massless and all values of $\langle a(x) \rangle$ are equally likely. Axion strings appear as topological defects. One must distinguish two cases: 1) inflation occurs with reheating temperature higher than the PQ transition temperature (equivalently, for our purposes, inflation does not occur at all) or 2) inflation occurs with reheating temperature less than the PQ transition temperature. In case 2 the axion field gets homogenized by inflation and the axion strings are 'blown' away.

When the temperature approaches the QCD scale, the potential $V(\theta)$ turns on and the axion acquires mass. There is a critical time, defined by $m_a(t_1)t_1 = 1$, when the axion field starts to oscillate in response to the turn-on of the axion mass. The corresponding temperature $T_1 \simeq 1 \text{ GeV}$ [29]. The initial amplitude of this oscillation corresponds to how far from zero the axion field is when the axion mass turns on. The axion field oscillations do not dissipate into other forms of energy and hence contribute to the cosmological energy density today [29]. This contribution is called of 'vacuum realignment'. It is further described below. Note that the vacuum realignment contribution may be accidentally suppressed in case 2 because the homogenized axion field happens to lie close to zero.

In case 1 the axion strings radiate axions [30,31]

from the time of the PQ transition till t_1 when the axion mass turns on. At t_1 each string becomes the boundary of N domain walls. If $N = 1$, the network of walls bounded by strings is unstable [32,33] and decays away. If $N > 1$ there is a domain wall problem [14] because axion domain walls end up dominating the energy density, resulting in a universe very different from the one observed today. There is a way to avoid this problem by introducing an interaction which slightly lowers one of the N vacua with respect to the others. In that case, the lowest vacuum takes over after some time and the domain walls disappear. There is little room in parameter space for that to happen and we will not consider this possibility further here. A detailed discussion is given in Ref. [34]. Henceforth, we assume $N = 1$.

In case 1 there are three contributions to the axion cosmological energy density. One contribution [30,31,35–39] is from axions that were radiated by axion strings before t_1 . A second contribution is from axions that were produced in the decay of walls bounded by strings after t_1 [36,40,41,34]. A third contribution is from vacuum realignment [29].

Let me briefly indicate how the vacuum alignment contribution is evaluated. Before time t_1 , the axion field did not oscillate even once. Soon after t_1 , the axion mass is assumed to change sufficiently slowly that the total number of axions in the oscillations of the axion field is an adiabatic invariant. The number density of axions at time t_1 is

$$n_a(t_1) \simeq \frac{1}{2} m_a(t_1) \langle a^2(t_1) \rangle \simeq \pi f_a^2 \frac{1}{t_1} \quad (7)$$

where f_a is the axion decay constant introduced earlier. In Eq. (7), we used the fact that the axion field $a(x)$ is approximately homogeneous on the horizon scale t_1 . Wiggles in $a(x)$ which entered the horizon long before t_1 have been red-shifted away [42]. We also used the fact that the initial departure of $a(x)$ from the nearest minimum is of order f_a . The axions of Eq. (7) are decoupled and non-relativistic. Assuming that the ratio of the axion number density to the entropy density

is constant from time t_1 till today, one finds [29]

$$\Omega_a \simeq \left(\frac{0.6 \cdot 10^{-5} \text{ eV}}{m_a} \right)^{\frac{7}{6}} \left(\frac{200 \text{ MeV}}{\Lambda_{QCD}} \right)^{\frac{3}{4}} h^{-2} \quad (8)$$

for the ratio of the axion energy density to the critical density for closing the universe. h is the present Hubble rate in units of 100 km/s.Mpc. The requirement that axions do not overclose the universe implies the constraint $m_a \gtrsim 6 \cdot 10^{-6} \text{ eV}$.

The contribution from axion string decay has been debated over the years. The main issue is the energy spectrum of axions radiated by axion strings. Battye and Shellard [37] have carried out computer simulations of bent strings (i.e. of wiggles on otherwise straight strings) and have concluded that the contribution from string decay is approximately ten times larger than that from vacuum realignment, implying a bound on the axion mass approximately then times more severe, say $m_a \gtrsim 6 \cdot 10^{-5} \text{ eV}$ instead of $m_a \gtrsim 6 \cdot 10^{-6} \text{ eV}$. My collaborators and I have done simulations of bent strings [36], of circular string loops [36,39] and non-circular string loops [39]. We conclude that the string decay contribution is of the same order of magnitude than that from vacuum realignment. Recently, Yamaguchi, Kawasaki and Yokoyama [38] have done computer simulations of a network of strings in an expanding universe, and concluded that the contribution from string decay is approximately three times that of vacuum realignment. The contribution from wall decay has been discussed in detail in ref. [34]. It is probably subdominant compared to the vacuum realignment and string decay contributions.

It should be emphasized that there are many sources of uncertainty in the cosmological axion energy density aside from the uncertainty about the contribution from string decay. The axion energy density may be diluted by the entropy release from heavy particles which decouple before the QCD epoch but decay afterwards [43], or by the entropy release associated with a first order QCD phase transition. On the other hand, if the QCD phase transition is first order [44], an abrupt change of the axion mass at the transition may increase Ω_a . If inflation occurs with reheat temperature less than T_{PQ} , there may be an accidental suppression of Ω_a because the homogenized ax-

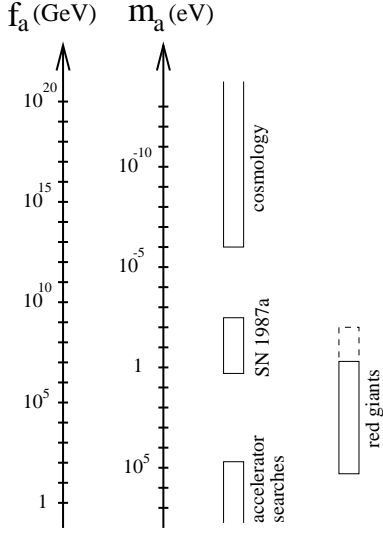


Figure 1. Ranges of axion mass m_a , or equivalently axion decay constant f_a , which have been ruled out by accelerator searches, the evolution of red giants, the supernova SN1987a, and finally the axion cosmological energy density.

ion field happens to lie close to a CP conserving minimum. Because the RHS of Eq. (7) is multiplied in this case by a factor of order the square of the initial vacuum misalignment angle $\frac{a(t_1)}{v}N$ which is randomly chosen between $-\pi$ and $+\pi$, the probability that Ω_a is suppressed by a factor x is of order \sqrt{x} . This rule cannot be extended to arbitrarily small x however because quantum mechanical fluctuations in the axion field during the epoch of inflation do not allow the suppression to be perfect [45].

The axions produced when the axion mass turns on during the QCD phase transition are cold dark matter (CDM) because they are non-relativistic from the moment of their first appearance at 1 GeV temperature. Studies of large scale structure formation support the view that the dominant fraction of dark matter is CDM. Any form of CDM necessarily contributes to galactic halos by falling into the gravitational wells of galaxies. Hence, there is excellent motivation to look for axions as constituent particles of our galactic halo.

Finally, let's mention that there is a particular kind of clumpiness [46,34] which affects axion dark matter if there is no inflation after the Peccei-Quinn phase transition. This is due to the fact that the dark matter axions are inhomogeneous with $\delta\rho/\rho \sim 1$ over the horizon scale at temperature $T_1 \simeq 1$ GeV, when they are produced at the start of the QCD phase-transition, combined with the fact that their velocities are so small that they do not erase these inhomogeneities by free-streaming before the time t_{eq} of equality between the matter and radiation energy densities when matter perturbations can start to grow. These particular inhomogeneities in the axion dark matter are in the non-linear regime immediately after time t_{eq} and thus form clumps, called 'axion mini-clusters' [46]. They have mass $M_{mc} \simeq 10^{-13}M_\odot$ and size $l_{mc} \simeq 10^{13}$ cm.

The various constraints on the axion, from accelerator searches, astrophysics and cosmology are summarized in Fig. 1.

4. The cavity detector of galactic halo axions

An electromagnetic cavity permeated by a strong static magnetic field can be used to detect galactic halo axions [47]. The relevant coupling is given in Eq. (5). Galactic halo axions have velocities β of order 10^{-3} and hence their energies $E_a = m_a + \frac{1}{2}m_a\beta^2$ have a spread of order 10^{-6} above the axion mass. When the frequency $\omega = 2\pi f$ of a cavity mode equals m_a , galactic halo axions convert resonantly into quanta of excitation (photons) of that cavity mode. The power from axion \rightarrow photon conversion on resonance is found to be [47,48]:

$$\begin{aligned}
 P &= \left(\frac{\alpha}{\pi} \frac{g_\gamma}{f_a} \right)^2 V B_0^2 \rho_a C \frac{1}{m_a} \text{Min}(Q_L, Q_a) \\
 &= 0.5 \cdot 10^{-26} \text{Watt} \left(\frac{V}{500 \text{ liter}} \right) \left(\frac{B_0}{7 \text{ Tesla}} \right)^2 \\
 &\quad \cdot C \left(\frac{g_\gamma}{0.36} \right)^2 \left(\frac{\rho_a}{\frac{1}{2} \cdot 10^{-24} \frac{\text{gr}}{\text{cm}^3}} \right) \\
 &\quad \cdot \left(\frac{m_a}{2\pi(\text{GHz})} \right) \text{Min}(Q_L, Q_a) \quad (9)
 \end{aligned}$$

where V is the volume of the cavity, B_0 is the magnetic field strength, Q_L is its loaded quality factor, $Q_a = 10^6$ is the ‘quality factor’ of the galactic halo axion signal (i.e. the ratio of their energy to their energy spread), ρ_a is the density of galactic halo axions on Earth, and C is a mode dependent form factor given by

$$C = \frac{\left| \int_V d^3x \vec{E}_\omega \cdot \vec{B}_0 \right|^2}{B_0^2 V \int_V d^3x \epsilon |\vec{E}_\omega|^2} \quad (10)$$

where $\vec{B}_0(\vec{x})$ is the static magnetic field, $\vec{E}_\omega(\vec{x})e^{i\omega t}$ is the oscillating electric field and ϵ is the dielectric constant.

Because the axion mass is only known in order of magnitude at best, the cavity must be tunable and a large range of frequencies must be explored seeking a signal. The cavity can be tuned by moving a dielectric rod or metal post inside it. Using Eq. (8), one finds the scanning rate to perform a search with signal to noise ratio s/n :

$$\begin{aligned} \frac{df}{dt} &= \frac{12\text{GHz}}{\text{year}} \left(\frac{4n}{s} \right)^2 \left(\frac{V}{500 \text{ liter}} \right)^2 \left(\frac{B_0}{7 \text{ Tesla}} \right)^4 \\ &\cdot C^2 \left(\frac{g_\gamma}{0.36} \right)^4 \left(\frac{\rho_a}{\frac{1}{2} \cdot 10^{-24} \frac{\text{gr}}{\text{cm}^3}} \right)^2 \left(\frac{3K}{T_n} \right)^2 \\ &\cdot \left(\frac{f}{\text{GHz}} \right)^2 \frac{Q_L}{Q_a}, \end{aligned} \quad (11)$$

where T_n is the sum of the physical temperature of the cavity plus the electronic noise temperature of the microwave receiver that detects the photons from $a \rightarrow \gamma$ conversion. Eq. (11) assumes that $Q_L < Q_a$ and that some strategies have been followed which optimize the search rate. The best quality factors attainable at present, using oxygen free copper, are of order 10^5 in the GHz range.

Eq. (11) shows that a galactic halo search with the required sensitivity is feasible with presently available technology, provided the form factor C can be kept at values of order one for a wide range of frequencies. For a cylindrical cavity and a homogeneous longitudinal magnetic field, $C = 0.69$ for the lowest TM mode. The form factors of the other modes are much smaller. The resonant frequency of the lowest TM mode of a cylindrical cavity is $f = 115 \text{ MHz} \left(\frac{1m}{R} \right)$ where R is the radius

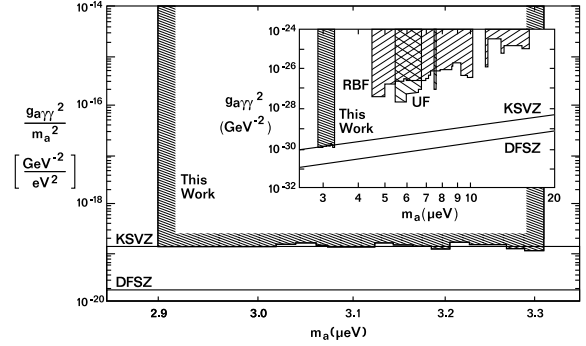


Figure 2. Axion couplings and masses excluded by the Large Scale US Dark Matter Axion Search at LLNL. Also shown are the KSVZ and DFSZ model predictions. Indicated on the insert are the regions excluded by the pilot experiments at Brookhaven National Laboratory (RBF) and the University of Florida (UF). All results are scaled to $\rho_a = 7.5 \cdot 10^{-25} \text{ g/cm}^3$.

of the cavity. Since $10^{-6} \text{ eV} = 2\pi (242 \text{ MHz})$, a large cylindrical cavity is convenient for searching the low frequency end of the range of interest. To extend the search to high frequencies without sacrifice in volume, one may power-combine many identical cavities which fill up the available volume inside a magnet’s bore [49,50]. This method allows one to maintain $C = 0(1)$ at high frequencies, albeit at the cost of increasing engineering complexity as the number of cavities increases.

Pilot experiments were carried out at Brookhaven National Laboratory [51] and at the University of Florida [52]. These experiments used relatively small magnets and hence the limits they placed on the local axion dark matter density are not severe. However they developed the various aspects of the cavity detection technique and demonstrated its feasibility.

Second generation experiments are presently under way at Lawrence Livermore National Laboratory (LLNL) [53] and at Kyoto University [54]. The LLNL experiment is similar in concept to the UF pilot experiment but uses a much larger magnet ($B_0^2 V = 12 T^2 m^3$). It is well engineered and runs with a near 100% duty cycle [55]. The results from its first year of running are reported in

ref. [56]. The exclusion plot is shown in Fig. 2. By definition, $g_{a\gamma\gamma} = \frac{\alpha}{\pi} \frac{g_{\gamma}}{f_a}$. The limits shown assume that the local halo density, estimated to be $7.5 \cdot 10^{-25} \text{ g/cm}^3$ [57], is entirely in axions. The experiment has ruled out the hypothesis that 100% of the local halo density is in KSVZ axions with mass in the range shown. Since then, the frequency range 500-800 MHz ($2.1 \leq m_a \leq 3.3 \mu\text{eV}$) has been searched but the results have not been published yet. Up till now, the experiment has used a single cavity with a variety of dielectric rods and metal posts. However, a four-cavity array will soon be used to search higher frequencies. Ultimately, the LLNL experiment will cover the mass range $1.3 \cdot 10^{-6} \text{ eV}$ to $13 \cdot 10^{-6} \text{ eV}$ at KSVZ sensitivity or better (see below).

A development project is under way to equip the LLNL detector with SQUID microwave receivers. These would replace the HEMT receivers presently in use. The HEMT receivers have noise temperature $T_n \sim 3 \text{ K}$ [58]. It appears that $T_n \lesssim 0.3 \text{ K}$ will be reached with the SQUIDS [59]. To take advantage of such low electronic noise temperatures, the experiment will have to be equipped with a dilution refrigerator. Also bucking coils must be installed to cancel the static magnetic field at the location of the SQUID. When this development project is completed, the LLNL detector will have sufficient sensitivity to detect DFSZ axions at even a fraction of the local halo density.

The Kyoto experiment exploits resonant $a \rightarrow \gamma$ conversion in a cavity permeated by a large static magnetic field, as do the other experiments, but uses a beam of Rydberg atoms to count the photons from $a \rightarrow \gamma$ conversion [54]. Single photon counting constitutes a dramatic improvement in microwave detection sensitivity. With HEMT amplifiers one needs to have thousands of $a \rightarrow \gamma$ conversions per second and integrate for about 100 sec to find a signal in the noise. With single photon counting, a few $a \rightarrow \gamma$ conversions suffice in principle. To build a beam of Rydberg atoms capable of single photon counting is a considerable achievement in itself. In addition, a dilution refrigerator is necessary to cool the cavity down to a temperature ($\sim 10 \text{ mK}$) where the thermal photon background is negligible. The projected

sensitivity of the Kyoto experiment is sufficient to detect DFSZ axions at even a fraction of the local halo density.

5. Other axion searches

There are a number of other techniques which have been used to search for very weakly coupled (so-called 'invisible') axions. Although these searches have not ruled out parameter space that is not also presently ruled out by the astrophysical limits described above, they do provide completely independent constraints.

5.1. Solar axion searches

The conversion of axions to photons in a magnetic field can be used to look for solar axions too [47,60–62]. The flux of solar axions on Earth is $\frac{7.4 \cdot 10^{11}}{\text{sec cm}^2} (\frac{g_{\gamma}}{0.36})^2 (\frac{m_a}{\text{eV}})^2$ from the Primakoff conversion of thermal photons in the sun [60]. The actual flux may be larger because other processes, such as Compton-like scattering, contribute if the axion has an appreciable coupling to the electron. At any rate the flux is huge compared to what can be produced by man-made processes on Earth and it is cost free. Solar axions have a broad spectrum of energies of order the temperature in the solar core, from one to a few keV.

Since the magnetic field is homogeneous on the length scale set by the axion de Broglie wavelength, the final photon is colinear with the initial axion. The photon and axion also have the same energy assuming the magnetic field is time-independent. The $a \rightarrow \gamma$ conversion probability is [47,60]

$$p = \frac{1}{4} \left(\frac{\alpha g_{\gamma}}{\pi f_a} \right)^2 (B_{0\perp} L F(q))^2 \quad (12)$$

if $B_{0\perp}(z) = B_{0\perp} b(z)$ is the magnetic field transverse to the direction of the colinear axion and photon, z is the coordinate along this direction, L is the depth over which the magnetic field extends and $F(q)$ is the form factor

$$F(q) = \frac{1}{L} \left| \int_0^L dz e^{iqz} b(z) \right| \quad (13)$$

where $q = k_{\gamma} - k_a = E_a - \sqrt{E_a^2 - m_a^2} \simeq \frac{m_a^2}{2E_a}$ is the momentum transfer. If the magnetic field is

homogeneous ($b = 1$), then

$$F(q) = \frac{2}{qL} \left| \sin \frac{qL}{2} \right|$$

$$\simeq 1 \quad \text{for } qL \ll 1 \quad . \quad (14)$$

For $qL \gg 1$, the conversion probability goes as $\sin^2(\frac{qL}{2})$ because the axion and photons oscillate into each other back and forth. The form factor $F(q)$ can be improved by filling the conversion region with a gas whose pressure is adjusted in such a way that the plasma frequency, which acts as an effective mass for the photon, equals the axion mass [60].

Multiplying the flux times the conversion probability, one obtains the event rate:

$$\frac{\text{events}}{\text{time}} \simeq \frac{200}{\text{day}} \frac{VL}{\text{meter}^4} F(q)^2 \left(\frac{B_{0\perp}}{8\text{Tesla}} \right)^2 \left(\frac{m_a}{\text{eV}} \right)^4 . \quad (15)$$

The final state photons are soft x-rays which may be detected with good efficiency. There are radioactive backgrounds to worry about however.

The above type of detector is usually referred to as an axion helioscope. If a signal is found due to axions or familons, the detector immediately becomes a marvelous new tool for the study of the solar interior. Experiments were carried out at Brookhaven National Lab. [61] and more recently at the University of Tokyo [62]. The BNL experiment used a stationary Isabelle dipole magnet whose aperture was directed towards the sun at sunset. The total exposure time was of order 15 minutes. The Tokyo experiment uses a superconducting dipole magnet mounted on an altazimuth which tracks the sun. The 95% confidence level upper limit [62] based on a few days of data taking is $g_{a\gamma\gamma} \equiv \alpha g_\gamma / \pi f_a \leq 6.0 \cdot 10^{-10} \text{GeV}^{-1}$ for $m_a \leq 0.03 \text{ eV}$.

F. Avignone and his collaborators [63] have exploited a different method to search for solar axions, namely the coherent Primakoff conversion of axions to photons in a crystal lattice. When the incident angle fulfills the Bragg condition for a given crystalline plane, the rate is enhanced. As the crystal detector turns with the Earth relative to the Sun's direction, a characteristic diurnal temporal pattern is produced. Using 1.94 kg.yr of data from a Ge detector in Sierra Grande, Argentina, the bound $g_{a\gamma\gamma} < 2.7 \cdot 10^{-9} \text{ GeV}^{-1}$

was obtained, independent of axion mass up to approximately 1 keV.

S. Moriyama [64] proposed looking for monochromatic axions emitted in the deexcitation of ^{54}Fe in the sun. ^{54}Fe has an M1 transition, between the first excited state and the ground state, with excitation energy 14.4 keV. The monochromatic axions can be resonantly absorbed by the same nucleus in the laboratory because the axions are Doppler broadened due to the thermal motion of the axion emitter in the Sun. An experiment of this type was carried out by M. Krčmar et al. [65].

5.2. Laser experiments

Eqs.(12,14) give the conversion probability in a static magnetic field of an axion to a photon of the same energy. The polarization of the photon is parallel to the component of the magnetic field transverse to the direction of motion. The inverse process, conversion of such a photon to an axion, occurs with the same probability p , of course. K. van Bibber et al. [66] proposed a 'shining light through walls' experiment in which a laser beam is passed through a long dipole magnet like those used for high-energy physics accelerators. In the field of the magnet, a few of the photons convert to axions. Another dipole magnet is set up in line with the first, behind a wall. Since the axions go through the wall unimpeded, this setup allows one to 'shine light through the wall.' An experiment of this type was carried out by the RBF collaboration [67]. Compared with a solar axion search, it has the advantage of greater control over experimental parameters. But the signal is much smaller because one pays twice the price of the very small axion-photon conversion rate.

Other types of laser experiments were proposed [68] in which one looks at the effect of the axion on the propagation of light through a magnetic field. If the photon beam is linearly polarized and the polarization direction is at an angle to the direction of the magnetic field, the plane of polarization turns because the component of light polarized parallel to the magnetic field gets depleted whereas the perpendicular component does not. There is an additional effect of birefringence because the component of light polarized parallel

to the magnetic field mixes with the axion and hence moves more slowly than in vacuo. Birefringence affects the ellipticity of the polarization as the light travels on. The birefringence associated with the axion is considerably smaller than that due to the box diagram in QED, i.e. an electron running in a loop with four external photon lines. The polarization rotation and the birefringence effects were searched for by the RBF collaboration [69], which boosted these effects by passing the laser beam hundreds of times in an optical cavity within the magnet. For more recent work, see ref. [70].

5.3. A telescope search

The axion decays to two photons at the rate:

$$\begin{aligned}\Gamma(a \rightarrow 2\gamma) &= g_\gamma^2 \frac{\alpha^2}{64\pi^3} \frac{m_a^3}{f_a^2} \\ &= \frac{g_\gamma^2}{6.8 \cdot 10^{24} \text{sec}} \left(\frac{m_a}{\text{eV}} \right)^5.\end{aligned}\quad (16)$$

For axions in the 10 eV mass range, the decay rate is comparable to the age of the universe. The dominant contribution to the cosmological energy density of such axions is thermal production [71]. The energy density in thermal axions is proportional to the axion mass and becomes equal to the critical energy density at a mass of order 100 eV, the exact value depending on how many particle species annihilate after the axion decoupled. One can search for relic axions by looking for the monochromatic photons from their decay. Such photons arrive to us from all directions but preferentially from large agglomerations of mass, such as clusters of galaxies. The relative width $\frac{\Delta\lambda}{\lambda}$ of the photon line from axion decay in galactic clusters is of order the virial velocity there, i.e. 10^{-2} . By subtracting the spectrum of light from a galactic cluster from that of the nightsky 'off cluster', one can subtract some of the background. The latter is dominated by lines in the spectrum of airglow [72]. A search of this type was carried out at Kitt Peak National Laboratory by the TSAR collaboration [72]. They placed the limit $g_{a\gamma\gamma} \leq 10^{-10} \text{GeV}^{-1}$ in the range $3 \leq m_a \leq 8 \text{ eV}$ (3100-8300 Å).

6. Macroscopic forces mediated by axions

J. Moody and F. Wilczek [73] analyzed the apparent deviations from the $1/r^2$ gravitational force law due to the exchange of virtual axions. The coupling of the axion to a spin 1/2 fermion f has the general form:

$$\mathcal{L}_{aff} = g_f \frac{m_f}{v} a \bar{f} (i\gamma_5 + \theta_f) f \quad (17)$$

when allowance is made for the fact that CP is violated in the electroweak interactions. g_f is of order one, whereas θ_f is of order $\bar{\theta}$, which is of order 10^{-17} [11,74] in the Standard Model. The second term in Eq. (17) produces a coupling of the axion field to the mass density of a macroscopic collection of non-relativistic fermions, whereas the first term produces a coupling of the axion field to the spin density of that macroscopic body. The first type of coupling was called 'monopole', the second 'dipole'. The axion mediated forces between two macroscopic bodies therefore fall into three categories: monopole-monopole, monopole-dipole, and dipole-dipole. The monopole-monopole force is suppressed by two powers of θ_f and therefore very small. The dipole-dipole force has a very large background from ordinary magnetic forces. This background can be suppressed by using superconducting shields but not well enough for the axion mediated contribution to be detected. The monopole-dipole is the least difficult to detect. It is suppressed by only one factor of θ_f and it can be modulated by rotating the spin polarized body. Experiments of this type have been carried out [75]. Unfortunately, their sensitivity is many orders of magnitude short of what is required to see the effect.

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